String Duality and
Novel Theories without Gravity[1]

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Abstract: We describe some of the novel 6d quantum field theories which have been discovered in studies of string duality. The role these theories (and their 4d descendants) may play in alleviating the vacuum degeneracy problem in string theory is reviewed. The DLCQ of these field theories is presented as one concrete way of formulating them, independent of string theory.

1 Introduction

Recent advances in string theory have led to the discovery of many new interacting theories without gravity. These theories are found by taking special limits of M-theory, in which many of the degrees of freedom decouple. In this talk we will:
I. Describe some examples of these new theories.
II. Review why it is important to fully understand these examples.
III. Propose a definition of these theories, in the light-cone frame, which is manifestly independent of M or string theory.

2 Examples

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[1]Based on a talk given at the "31st International Symposium Ahrenshoop on the Theory of Elementary Particles," Buckow, Germany, September 2-6 1997

2.1 Theories with $(2,0)$ Supersymmetry

The first (and simplest) examples were found by Witten [1], in studying type IIB string theory on $K3$. He considers the situation where the $K3$ develops an $A - D - E$ singularity. In the IIA theory, one finds extra massless gauge bosons in these circumstances. These extra vectors of the (1,1) supersymmetry are required by string-string duality, and arise from D2 branes which wrap the collapsing 2-cycles and become massless in the singular limit.

In the IIB theory, there is a chiral (2,0) supergravity in six dimensions. The only massless multiplet of the (2,0) supersymmetry (other than the gravity multiplet) is the tensor multiplet, which consists of 5 scalars, some chiral fermions, and a self-dual two form $B_{\mu\nu}$ which satisfies

$$dB = *dB \tag{1}$$

The (2,0) supergravity *requires* the presence of precisely 21 tensor multiplets for anomaly freedom. Therefore, it is hard to envision a scenario where one finds extra massless particles at the singular point in moduli space. However, further compactification on an S^1 yields a theory related to the IIA theory by T-duality, so one must find (after S^1 compactification) gauge bosons of the $A - D - E$ gauge group. What is their IIB origin?

Further compactify the IIA and IIB theories on circles with radii $R_{A,B}$. Then T-duality relates the theories with $R_A = \frac{1}{R_B}$ (we are temporarily setting the string scale α' to one for simplicity). The relation between the six-dimensional string couplings $\lambda_{A,B}$ is

$$\frac{1}{\lambda_A} = \frac{R_B}{\lambda_B} \tag{2}$$

If we consider a point in IIA moduli space a distance ϵ from the singular point, then there are W-bosons coming from wrapped D2 branes whose masses go like

$$M_W = \frac{\epsilon}{\lambda_A} \tag{3}$$

So in type IIB, the mass is

$$M_W = \frac{\epsilon R_B}{\lambda_B} \tag{4}$$

This looks like the mass of a *string* wrapped around the S^1 in the IIB theory! But this string is not the critical type IIB string; from equation (4) it must have a tension

$$T = \frac{\epsilon}{\lambda_B} \tag{5}$$

Of course, this string comes from a D3 brane wrapped around a collapsing sphere in the $K3$ of area $\simeq \epsilon$.

For very small ϵ, $T << \frac{1}{\alpha'}$. So we get an $A - D - E$ series of quantum theories in six dimensions which contain light string solitons. Because the noncritical strings are very light compared to the fundamental string scale, one can decouple gravity. Then, it is believed that one is left with an interacting quantum field theory in six

dimensions. As $\epsilon \to 0$, one approaches a nontrivial fixed point of the renormalization group.

In six-dimensions, strings are dual to strings. The particular light strings in question are self-dual (the $H = dB$ they produce is self-dual as in equation (1)), so the "coupling" of these quantum theories is fixed and of order one. In other words, there is no coupling constant which can serve as an expansion parameter.

By using ALE spaces instead of $K3$, one can find such interacting theories for each A_k or D_k singularity. For the A_k theories, there is another simple description due to Strominger and Townsend [2, 3]. For instance, consider two parallel M5 branes in eleven-dimensional flat spacetime. There are membranes which can end on the M5 branes, yielding a noncritical string on the fivebrane worldvolume with tension proportional to the separation.
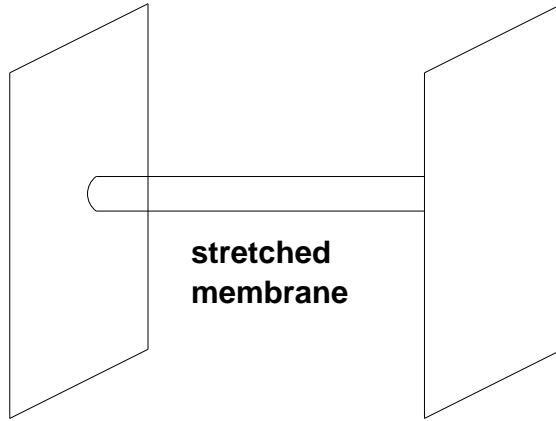

Figure 1: A membrane stretching between two M5 branes.

Each fivebrane has a tensor multiplet on its worldvolume (the five scalar components parametrize the transverse position of the fivebrane in eleven dimensions). If we denote the two five-tuples of scalars by $\vec{\phi}_{1,2}$, then the VEVs $\langle \vec{\phi}_{1,2} \rangle$ label different vacua of an effective six-dimensional theory. When $\vec{\phi}_1 \to \vec{\phi}_2$, the noncritical strings become tensionless and we find another description of the A_1 fixed point above.

More precisely, if we say the fivebranes are separated by a distance L, the limit one wishes it to take is

$$M_{pl} \to \infty, \quad L \to 0, \quad T = LM_{pl}^3 \; fixed \tag{6}$$

In this limit, gravity decouples but the noncritical strings stay light, yielding an interacting theory without gravity. The obvious generalization with k fivebranes yields the A_{k-1} (2,0) fixed point, with moduli space

$$\mathcal{M}_k = \frac{R^{5k}}{S_k} \tag{7}$$

3

given by the positions of the parallel fivebranes, mod permutations. At generic points on $\mathcal{M}_k$, the low energy theory has k tensor multiplets.

2.2 Theories with $(1,0)$ Supersymmetry

New interacting 6d theories with (1,0) supersymmetry have also been discovered [4, 5]. Perhaps the simplest example is the following. Consider Horava and Witten's description of the $E_8 \times E_8$ heterotic string as M-theory on S^1/Z_2 [6] . The length of the interval is related to the heterotic g_s, while E_8 gauge fields live on each of the two "end of the world" ninebranes.

We can consider a fivebrane at some point on the interval. Its position in the S^1/Z_2 is parametrized by the real scalar ϕ in a (1,0) tensor multiplet, while its other transverse positions are scalars in a (1,0) hypermultiplet. Since it is a scalar in six dimensions, ϕ naturally has dimension two; we will say the two E_8 walls are located at $\phi = 0$ and $\phi = \frac{1}{\alpha'}$. Then, one has noncritical strings on the fivebrane world volume with tensions

$$T_1 = \phi, \quad T_2 = (\frac{1}{\alpha'} - \phi) \tag{8}$$

coming from membranes with one end on the fivebrane and one end on the nine-branes.

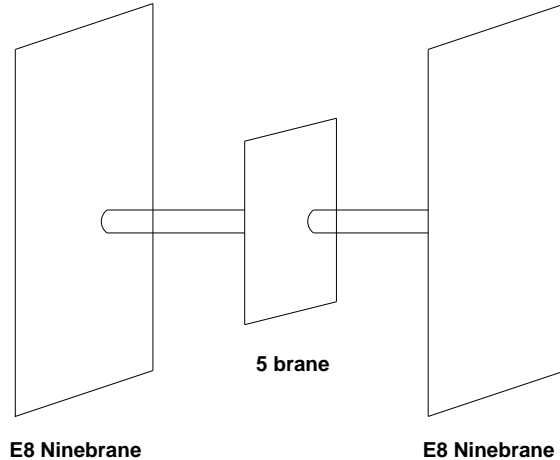

Figure 2: A fivebrane between two ninebranes and two 5-9 membranes.

As $\phi \to 0$ or $\phi \to \frac{1}{\alpha'}$, one again finds that the lightest degrees of freedom in the theory are solitonic self-dual noncritical strings. The theory on the fivebrane needs to also have an E_8 global symmetry, to couple consistently to the E_8 gauge fields on the end of the world. Therefore, one concludes that when the fivebrane hits the ninebrane, one finds a nontrivial (1,0) supersymmetric RG fixed point, with E_8 global symmetry.

For both the (2,0) A_k theories and the E_8 theory in 6d, there is no known UV free Lagrangian which flows, in the IR, to the fixed point of interest. Therefore, it is

of intrinsic interest to find a definition of these quantum field theories which is independent of string or M-theory. We will propose such a definition in §4. Before doing that, it seems proper to provide some motivations for the study of these theories.

3   Why are these theories of interest?

There are at least three motivations for studying these theories:

a) These are the first examples of nontrivial fixed point quantum field theories above four dimensions. For instance, if one does a very naive analysis of gauge field theory in d dimensions

$$L \;=\; \int d^d x \; \frac{1}{g^2} F_{\mu\nu}^2 + \cdots \tag{9}$$

one finds that $[g^2] = 4 - d$ (in mass units), so the theory is *infrared free* for $d > 4$. Hence, the theories of §2 are of intrinsic interest as a new class of interacting quantum field theories.

b) These theories play a crucial role in the study of M(atrix) theory compactifications. In M(atrix) theory, one starts with the maximally supersymmetric D0 brane quantum mechanics, with N zero branes giving the DLCQ in a sector with light-like momentum N and the $N \to \infty$ limit yielding the 11d uncompactified theory [8, 9]. To study compactifications on a transverse T^n, one then T-dualizes the $U(N)$ D0 brane quantum mechanics to obtain a description of M-theory on T^n as the $n+1$ dimensional $U(N)$ Super Yang-Mills theory compactified on the dual torus, $\tilde{T}^n$.

An obvious problem with this approach is that for $n > 3$, the Super Yang-Mills is ill-defined at short distances (it is not renormalizable). Let us consider the first such case: M-theory on T^4. The U-duality group in this case is $SL(5, Z)$. This suggests that perhaps the M(atrix) definition involves some 5+1 dimensional QFT compactified on a $\tilde{T}^5$, geometrizing the $SL(5, Z)$ U-duality group as the modular group of the torus. The unique candidate which is well-defined (and has the correct supersymmetry) is the A_N (2,0) quantum field theory of §2 [10, 11]. But, how does this prescription relate to our expectation that the theory should be 4+1 $U(N)$ SYM?

The 5d SYM theory has a conserved $U(1)$ current

$$j = *(F \wedge F) \tag{10}$$

We can identify j with the Kaluza-Klein $U(1)$ symmetry of the (2,0) theory compactified on an S^1 of radius $\tilde{L}_5$, if we say

$$2\pi\tilde{L}_5 \;=\; \frac{g^2}{2\pi} \tag{11}$$

The 5d theory has particles which are 4d instantons, and whose action is given by $4\pi^2 \frac{n}{g^2}$ for the "n-instanton" particle. These in turn can be interpreted as Kaluza-Klein modes coming from the (2,0) theory with a momentum $p_5 = \frac{n}{\tilde{L}_5}$ around the hidden "extra" circle which promotes $\tilde{T}^4$ to $\tilde{T}^5$. In this way, one ends up with the

prescription that M-theory on T^4 is defined by the (2,0) theory on $\tilde{T}^5$ [10, 11]. This makes the SL(5,Z) U-duality manifest.

c) These novel interacting theories play a crucial role in the unification of M-theory vacua. Consider, for instance, the heterotic $E_8 \times E_8$ theory compactified on $K3$. There is a Bianchi identity for the three-form field strength H which looks like

$$dH = Tr(R \wedge R) - Tr(F \wedge F) \tag{12}$$

where R is the curvature and F is the Yang-Mills field strength. Integrating (12) over the $K3$, we find that there should be $n_{1,2}$ Yang-Mills instantons in the two E_8s, with

$$n_1 + n_2 = 24 . \tag{13}$$

It is then natural to ask: How are vacua with different choices of $n_{1,2}$ connected to each other?

Consider the n_1 instantons in one E_8 wall. Instanton moduli space has singularities, including points where a single instanton shrinks to "zero size." In the heterotic theory, this small instanton can now be represented as a fivebrane sitting at the E_8 wall. But, now there is a new branch in the moduli space of vacua - in addition to re-expanding into a large E_8 instanton, the fivebrane can move off into the S^1/Z_2 interval! In the process, one loses 29 hypermultiplet moduli (the moduli of one E_8 instanton) and gains a single tensor multiplet (the real scalar parametrizes the position of the fivebrane in the interval). Hence, one is left with $n_1 - 1$ instantons on the E_8 wall. By moving across the interval and entering the other wall as an instanton, the fivebrane can effect a transition from a vacuum with instanton distribution (n_1, n_2) to a vacuum with instanton distribution (n_1-1, n_2+1). In this way, the perturbative heterotic vacua with different numbers of instantons in the two E_8s are all connected [4, 5]. More generally, one can modify equation (13) to read $n_1 + n_2 + n_5 = 24$, where now n_5 is the number of five-branes in the interval [7].

We have glossed over an important point here: The (1,0) tensor multiplet (on the fivebrane worldvolume) contains a self-dual tensor $B_{\mu\nu}^+$. No conventional mass term is possible for the tensor, since there is no $B_{\mu\nu}^-$ that $B_{\mu\nu}^+$ can pair up with. So, how can transitions changing the number of tensor multiplets ever occur?

When the transition occurs, we are precisely in the situation described in §2.2, where there is an interacting (1,0) superconformal field theory with E_8 global symmetry. There is no weakly coupled description of this fixed point, and a phase transition can occur there. By going through this nontrivial fixed point, it is possible to connect the two branches with different numbers of tensor multiplets. So, the novel theories of §2 are of apparent use in unifying 6d (0,1) supersymmetric vacua.

In fact, related theories also seem to play an important role in connecting 4d $N = 1$ vacua. For instance, one can compactify the $E_8 \times E_8$ heterotic string on a Calabi-Yau threefold M which is a $K3$ fibration. In many examples, one finds a low-energy theory with *chiral* gauge representations. For instance, if one has embedded an $SU(3)$ bundle V with $c_1(V) = 0$ in one of the E_8s, the unbroken subgroup of E_8 is E_6. The matter fields come in the **27** and $\overline{\textbf{27}}$ representations, and one finds for

the net number $N = |\#\mathbf{27} - \#\overline{\mathbf{27}}|$ of generations:

$$N = \frac{1}{2}| \int_M c_3(V)| \ . \tag{14}$$

Among the singular loci in the moduli space of vacua, there are places where V develops a curve of singularities which corresponds to a small instanton in the generic $K3$ fiber of M. One can represent this small instanton as a fivebrane wrapping the base of the fibration. In certain cases, there is a new branch of the moduli space where the wrapped fivebrane can move away from the E_8 wall into the S^1/Z_2. It was argued in [12] that in many cases this changes the net number of generations of the E_6 gauge theory remaining on the wall. Related phenomena were discussed in [13, 14]. So, phase transitions through close relatives of the E_8 fixed point in six dimensions can also connect up 4d string vacua with different net numbers of generations.

4 A Proposed M(atrix) Description

In §3, we have seen several interesting applications of the new interacting 6d field theories. However, for both the (2,0) supersymmetric theories and the (1,0) theories with E_8 global symmetry, there is no obvious definition of the theory that doesn't involve an embedding in M-theory. This is a very $un - economical$ way of defining a quantum field theory – one starts with far too many degrees of freedom, and must decouple most of them from the quantum field theory of interest.

An alternative way of describing the (2,0) theories was proposed in [15, 16], and extended to the (1,0) theories in [17, 18]. We will discuss the simplest case – the A_{k-1} (2,0) theory, i.e. the theory of k coincident M 5-branes. We know several suggestive facts about this theory:
• If we compactify the 6d theory on a circle with radius R, it produces a 5d $U(k)$ Super Yang-Mills theory with coupling $g_5^2 = R$.
• The Kaluza-Klein particles (with $p_5 = \frac{1}{R}$) are "instantons" of the $U(k)$ gauge theory (i.e., 4d Yang-Mills instantons which look like particles in 5d).
These facts suggest that, in analogy with the M(atrix) approach to M-theory [8], we should search for a light-cone quantization of the A_{k-1} (2,0) theory.

4.1 DLCQ of (2,0) A_{k-1} Theory

Let us take our 5+1 dimensions to be parametrized by $X^0, \cdots, X^5$. In normal light-cone quantization, one defines $X^\pm = X^0 \pm X^1$ and gives initial conditions on a surface of fixed X^+. Then, one evolves forward in light-cone "time" using the Hamiltonian $H = P_+$. The modes of quantum fields with $P_- < 0$ are canonical conjugates of modes of $P_- > 0$, so we can choose the vacuum to be annihilated by the $P_- < 0$ modes. The $P_- = 0$ modes are not dynamical (but can give rise to subtleties, which will be mentioned below).

In *discrete* light-cone quantization (or DLCQ) [9], one in addition compactifies the light-like direction

$$X^- \simeq X^- + 2\pi R \tag{15}$$

Then, P_- is quantized in units of $1/R$. For finite N, the DLCQ Fock space is very simple, since there are a finite number of modes. However integrating out the zero momentum modes can, in the DLCQ of *some* theories, lead to complicated interactions [19]. The decompactification limit (where one recovers Lorentz invariance) is taken by going to large R at fixed P_-, which is equivalent to going to large N.

Following [20], one may find it fruitful to view the compactification of X^- as the limit of a space-like compactification

$$X^- \simeq X^- + 2\pi R, \quad X^+ \simeq X^+ + \frac{R_s^2}{R} \tag{16}$$

where $R_s \to 0$. For finite R_s, one can boost this to a spatial compactification

$$X^1 \sim X^1 + R_s, \quad P_- = \frac{N}{R} \to P_1 = \frac{N}{R_s} \tag{17}$$

Then, our interest is in describing modes of momentum P_1 as $R_s \to 0$. In many cases, this can be rather complicated. But for the (2,0) superconformal theories, we get a very weakly coupled ($g^2 = R_s$) $U(k)$ Super Yang-Mills theory in 4+1 dimensions, with N "instantons" (carrying charge under $J = *(F \wedge F)$). If we want these to have finite energy in the *original* reference frame, they must have very small velocities. So for $R_s \to 0$, it seems that we should get a quantum mechanical sigma model on the moduli space of N $U(k)$ instantons. The space-time supersymmetry implies that this sigma model should have 8 supercharges.

We will now give a more direct derivation of the relevant quantum mechanics from M(atrix) theory.

4.2   Derivation from M(atrix) Theory

Following Berkooz and Douglas [21], we know that the background of k longitudinal 5-branes in M(atrix) theory can be represented by studying the theory of N zero branes in the background of k D4 branes in Type IIA string theory. The presence of the D4 branes (and the consequent 0-4 strings) break the supersymmetry of the quantum mechanics to $N = 8$.

The resulting quantum mechanics is in fact the dimensional reduction of a 6d (0,1) supersymmetric system, which is a $U(N)$ gauge theory with k fundamental hypermultiplets and an additional adjoint hyper. The quantum mechanics has a $U(N)$ gauge symmetry and an $SO(4)_\parallel \times SO(5)_\perp \times U(k)$ global symmetry. The bosonic fields are $X_\perp$, $X_\parallel$ and q with charges $(\mathbf{N^2}, \mathbf{1}, \mathbf{5}, \mathbf{1})$, $(\mathbf{N^2}, (\mathbf{2}, \mathbf{2}), \mathbf{1}, \mathbf{1})$ and $(\mathbf{N}, (\mathbf{2}, \mathbf{1}), \mathbf{1}, \mathbf{k})$ under the gauge and global symmetries. Roughly speaking, $X_\perp$ characterizes the positions of the 0 branes transverse to the 4-branes, while $X_\parallel$ characterizes their positions in the directions along the 4-branes.

The moduli space of vacua has various branches, but two are of particular interest to us [2]:

1) The Coulomb branch

On this branch, $X_\perp, X_\parallel \neq 0$ while $q = 0$. This is the branch where the 0 branes are moving around in spacetime away from the D4 branes, as depicted in Figure 3 below.
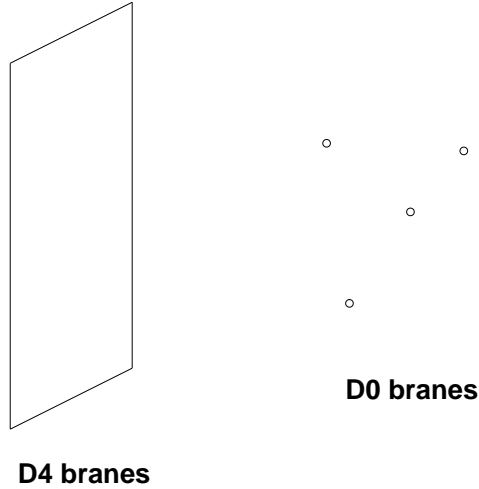

D0 branes

D4 branes
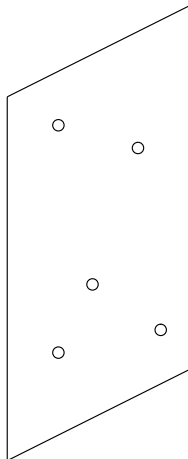
Figure 3: A picture of a point on the Coulomb branch.

2) The Higgs branch

On this branch $X_\perp = 0$ while $X_\parallel, q \neq 0$. The D0 branes are inside of the D4 branes, as shown below in Figure 4. We will denote the Higgs branch moduli space for given N and k by $\mathcal{M}_{N,k}$.

So roughly speaking, the Higgs branch is concerned with the physics on the fivebranes while the Coulomb branch captures the physics away from the fivebranes (e.g. 11d supergravity). This leads us to believe that the quantum mechanics on the Higgs branch will offer a M(atrix) description of the interacting field theories discussed in §2.1. More precisely, if we want to decouple gravity from the physics on the fivebranes, we need to take the limit $M_{pl} \to \infty$. Then in particular, the coupling in the quantum mechanics g_{QM} which is related to R and M_{pl} by $g_{QM}^2 = R^3 M_{pl}^6$ also satisfies $g_{QM} \to \infty$. This is the infrared limit of the quantum mechanics.

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[2]Of course strictly speaking the quantum mechanics has no moduli space of vacua, but as usual in discussions of M(atrix) theory we can imagine a moduli space in the Born-Oppenheimer approximation.

D0 branes inside D4 branes

Figure 4: A picture of a point on the Higgs branch.

The surprising fact is that in this strong coupling limit, many simplifications occur:

1) For $g_{QM} \to \infty$, the Higgs branch physics decouples from the Coulomb branch! This is because the masses of the massive W bosons go off to infinity with g_{QM}, and there is a tube of infinite length in the Coulomb branch as one approaches $X_\perp \to 0$. So, we are left with quantum mechanics on the Higgs branch $\mathcal{M}_{N,k}$.

2) By the ADHM construction, $\mathcal{M}_{N,k}$ is actually the moduli space of N instantons in $U(k)$ gauge theory [22]! This is in accord with the fact that D0 branes are expected to behave like instantons in D4 branes [23].

3) By theorems about the absence of couplings between vector and hypermultiplets, the Higgs branch does *not* receive quantum corrections. Therefore, even the strong coupling $g_{QM} \to \infty$ limit does not correct the sigma model on $\mathcal{M}_{N,k}$.

We conclude that the A_{k-1} (2,0) theory has a description in terms of the quantum mechanics on the moduli space of N $U(k)$ instantons in the $N \to \infty$ limit. One can similarly derive a M(atrix) description for the E_8 theories with (1,0) supersymmetry [17, 18]. We will not review that here.

5   Conclusions

I have tried to emphasize in this talk the fruitful interaction that has occurred in the past year or two between three different research directions. The new interacting theories in $d \geq 4$ seem to have important uses in both the quest for a nonperturbative definition of (compactified) M-theory, and in the search for resolutions to the vacuum degeneracy problem.

Several advances have been made in areas very closely related to my talk since the conference in Buckow. I summarize some recent developments and future directions here:

• One can try to use similar M(atrix) formulations to study other interesting field theories, for instance familiar 4d field theories [24].

• One can use the quantum mechanical formulation of the 6d theories to try and compute quantities of interest in these conformal field theories, e.g. correlation functions of various local operators [25].

• One can try to use the M(atrix) descriptions of the "little string theories" of [26] to compute interesting properties of these novel theories without gravity.

• One can investigate M-theory compactification on six and higher dimensional manifolds in the M(atrix) formulation. There are problems with obtaining a simple M(atrix) description of T^6 compactifications (as discussed in e.g. [20]), while the situation seems to be better for Calabi-Yau compactifications [27]. Successful definitions of the M(atrix) compactifications will involve new theories without gravity.

• One can further pursue the study of interesting phase transitions in 4d $N = 1$ string vacua [12, 13, 14] by using new constructions (utilizing D-branes or F-theory) to find simple examples.

Acknowledgements

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